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Status of weak-scale supersymmetry

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Abstract : This article includes discussions on :
(i) Standard Model and motivation for supersymmetry,
(ii) Supersymmetry and MSSM,
(iii) CMSSM and the mass spectra of sparticles,
(iv) Experimental constraints on CMSSM parameters and
(v) Conclusions

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1. Standard Model and motivation for supersymmetry

Supersymmetry, as a generalized spacetime invariance under which fermions and bosons transform into each other, is undoubtedly a beautiful idea. But why should particle physicists look for it—especially at or slightly above the weak scale? The answer is that softly broken supersymmetry with an intra-supermultiplet mass breaking ≤ 0 (TeV) can cure the Standard Model (SM) of particle physics of a serious theoretical deficiency, *viz.* the radiative instability of the Higgs mass.

Before I elaborate on the last point, let me first briefly review the impressive experimental successes [1] that SM has had—if only to underscore the absence of any phenomenological need at present to go beyond it. Table 1 below contains a "pull-plot". Various measurables (on the Z-peak) of the standard electroweak theory have been listed in the first column. The second column contains the measured values (combining SLD and LEP numbers) with 1σ errors. The "pull", defined as the deviation (with sign) of the central value from the theoretical prediction divided by the 1σ error in the measurement, is given in the third column. The fourth column displays the same information geometrically in terms of horizontal bars drawn in units of σ.

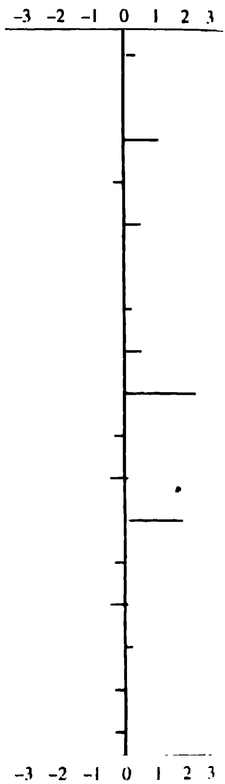
In Table 1 there are seventeen data items and one fitting parameter, namely the Higgs mass m_H, χ^2/(degree of freedom) being 18.5/15 in the fit. The best fit for the latter is [1]

$$m_H = 149^{+148}_{-78}\ \text{GeV},$$

to be contrasted with the latest result $m_H > 70.5$ GeV from direct search experiments at LEP. One would readily agree that the data represent an outstanding success of the EW

Table 1. Pull-plot for electroweak measurables on the Z-peak

M_Z [GeV]	91.1863 ± 0.0020	.17
Γ_Z [GeV]	2.4946 ± 0.0027	.05
σ^0_{hadr} [nb]	41.508 ± 0.056	.97
R_e	20.754 ± 0.057	-.22
R_μ	20.796 ± 0.040	.73
R_τ	20.814 ± 0.055	.00
$A^{0,e}_{fb}$	0.0160 ± 0.0024	.32
$A^{0,\mu}_{fb}$	0.0162 ± 0.0013	.74
$A^{0,\tau}_{fb}$	0.0201 ± 0.0018	2.70
A_τ	0.1401 ± 0.0067	-.37
A_e	0.1382 ± 0.0076	-.57
R_b	0.2179 ± 0.0012	1.70
R_c	0.1715 ± 0.0056	-.14
$A^{0,b}_{fb}$	0.0979 ± 0.0023	-.87
$A^{0,c}_{fb}$	0.0733 ± 0.0049	.43
$\sin^2\bar{\theta}^{e}_{w}$	0.2320 ± 0.0010	-.09
$1/\alpha$	128.894 ± 0.090	-.25



electroweak sector of SM, though mild doubts can be entertained regarding the $\tau^+\tau^-$ forward-backward asymmetry at the Z and the Z-decay branching fraction into $b\bar{b}$. Turning to QCD [2], Figure 1 shows a "best fit" plot of the QCD fine structure coupling α_S evolving *via* the renormalization group equation as a function of the energy scale Q. Given the large number of different determinations at different scales, one would call the agreement quite impressive.

In the light of such an outstanding experimental success, any theoretical motivation for going beyond SM needs to be compelling. Such a motivation was indeed put forth by 't Hooft in 1980 by showing the radiative instability of the Higgs mass : a feature of SM known as the naturalness problem. Already, at the 1-loop level, the presence of a quadratic divergence in the summed diagrams of Figure 2 implies the following fact. If there are

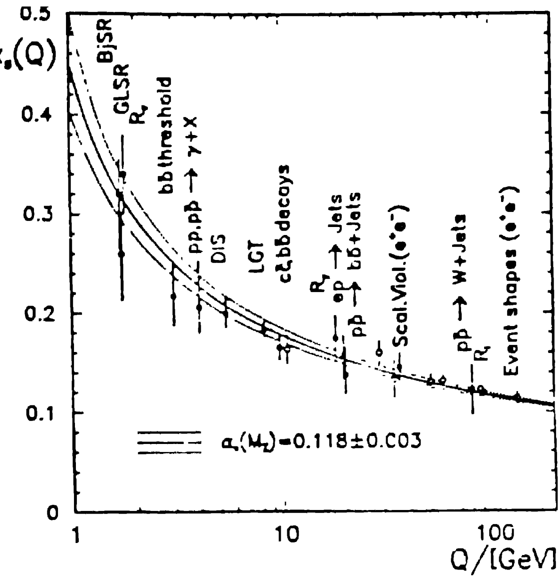


Figure 1. Evolution of α_S with Q. QCD theory vs experiment.

unknown superheavy fields at some high scale M (*e.g.* Planck scale M_{Pl}), SM has to be viewed as a residual theory at low energies after these superheavy fields have been integrated out. However, the latter procedure makes the finite mass of the electroweak Higgs shift quadratically to that high scale M. An unnatural amount of fine tuning is needed

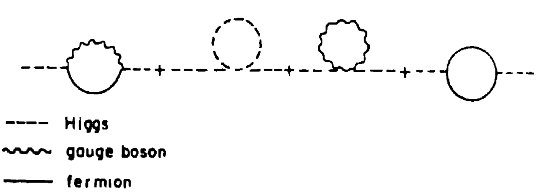


Figure 2. 1-loop contributions to the Higgs mass.

order by radiative order between the Higgs mass and self-coupling parameters in the Lagrangian to keep m_H within an electroweak range. Generally, one can represent the effect of integrating the superheavy fields as :

$$\mathcal{L}\left(\phi_{\text{light}}, \phi_{\text{heavy}}\right) \to \mathcal{L}^{\text{eff}}\left(\phi_{\text{light}}\right) = \mathcal{L}\left(\phi_{\text{light}}\right) + \sum c_n O_n.(M)^{-d_n+4}. \quad (1)$$

where $\mathcal{L}^{\text{eff}}\left(\phi_{\text{light}}\right)$ remains the effective Lagrangian for the residual theory. In (1) d_n is the scale dimension of the operator O_n in the operator product expansion of the RHS. The problem with the Higgs mass term is that the leading value of d_n in the RHS of (1) is 2 so that M comes in as M^2 and the high scale contribution evidently does not become smaller as M increases. The fine tuning would be to adjust the corresponding coefficient c_n to zero.

There have traditionally been two types of suggestions for the way out : (1) the strong coupling (new structure) option and (2) the weak coupling (new symmetry) option. (1) is presently disfavoured by precision tests in which SM has performed very well. In particular, the electroweak oblique parameters S, T, U (which vanish for SM) are now experimentally know to be [1] $-0.04^{+0.17}_{-0.11}$, $-0.18^{+0.17}_{0.08}$ and 0.07 ± 0.42 respectively. These are generically expected to be $O(1)$ in option (1). A similar negative conclusion regarding this option follows from rather strong upper limits which exist on any flavor-changing weak neutral current. In option (2) supersymmetry is perceived to be the new desired symmetry. Here quadratic divergences of fermionic and bosonic loops cancel with opposite signs and any radiative shift in the Higgs mass squared gets controlled by the squared mass-difference between particles and their superpartner sparticles within the same supermultiplet. So long as the latter is ≤ 0 (TeV^2), there is no problem. Referring back to (1), the coefficient c_n of O_n with $d_n = 2$ is naturally made to vanish by supersymmetry.

The power of supersymmetry can best be understood from a simple toy model : scalar electrodynamics. The mass of the scalar field in this theory is unprotected against large radiative corrections by any symmetry and suffers from the naturalness problem owing to a quadratic divergence present already at the 1-loop level. The fermion mass in spinor electrodynamics, on the other hand, is protected by chiral symmetry and does not have this problem—the corresponding loop divergences being logarithmic In supersymmetric quantum electrodynamics (SQED), the mass of the scalar is equal to the mass of its partner fermion by supersymmetry and hence gets protected. Thus SQED, unlike scalar QED, is a natural theory. Moreover, this feature persists even with supersymmetry breaking so long as the latter is done by soft terms (*i.e.* of scale dimensions less than four)

2. Supersymmetry and MSSM

The supersymmetry idea, originally due to Golfand and Likhtmann [3], was developed further by Akulov and Volkov and more specifically in the context of quantum field theory by Wess and Zumino as well by Salam and Strathdee. It postulates the existence of particle-sparticle supermultiplets with the superpartners differing in spin by 1/2 unit. Thus a supersymmetric theory contains supermultiplets with spins 0 and 1/2 (*e.g.* quarks and squarks or electrons and selectrons or Higgs bosons and higgsinos) as well as those with spins 1 and 1/2 (*e.g.* photon and photino or gluons and gluinos or W's and winos or Z and zino *etc.*). Neutral higgsinos mix with the zino and the photino into four physical neutralinos, while charged higgsinos and winos mix into two pairs of physical charginos. The new particles (called sparticles) become necessary since established quantum numbers forbid one to make supermultiplets out of the known fermions and bosons. In the local version of supersymmetry there is also the supermultiplet comprising the spin 2 graviton and the spin 3/2 gravitino.

The zoo of sparticles, as well as the symbols for themselves and their superfields, appears in Table 2 below while Figure 3 graphically shows how different particles and sparticles are denoted by characteristic lines in Feynman diagrams. These sparticles should, in general, have masses characterized by the intra-supermultiplet splitting scale M_s where $M_w < M_s \leq 0$ (TeV). In particular, if all extra particles—necessitated by a supersymmetric extension of the Standard Model—are heavier than 200 GeV, supersymmetry will decouple [4] at presently available energies. Residually left will be the Standard Model in its pristine form with a somewhat light Higgs particles.

Table 2. Zoo of sparticles.

Name	Symbol
sleptons	$\tilde{l}_{L,R}$
(selectron, smuon, stau)	$\left(\tilde{e}_{L,R}, \tilde{\mu}_{L,R}, \tilde{\tau}_{L,R}\right)$
squarks	$\tilde{q}_{L,R}$
(*s*-up, *s*-down, *s*-charm,	$\left(\tilde{u}_{L,R}, \tilde{d}_{L,R}, \tilde{s}_{L,R},\right.$
s-strange, stop, sbottom)	$\left.\tilde{s}_{L,R}, \tilde{t}_{L,R}, \tilde{b}_{L,R}\right)$
gluino	$\tilde{g}$
charginos	$\tilde{\chi}^{\pm}_{1,2}$
neutralinos	$\tilde{\chi}^{0}_{1,2,3,4}$
gravitino	$\tilde{G}$

The minimal supersymmetric extension of the Standard Model, *i.e.* the one with the minimum number of extra particles is called the Minimal Supersymmetric Standard Model

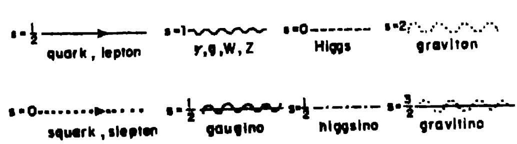


Figure 3. Legend for describing particles and sparticles in Feynman diagrams.

MSSM [5]. Its spectrum consists of the particles of SM—with a minimally extended Higgs sector—and their partner sparticles. For particles, the only new feature, as already mentioned, is that—in place of one—there are five physical Higgs scalars (a charged pair $H^{\pm}$, two *CP*-even neutrals—the lighter h and the heavier H—as well as one *CP*-odd neutral A) emerging from two Higgs doublets which occur here instead of one as in SM. The ratio = (VEV of the neutral Higgs field which couples to up-type fermions) ÷ (VEV of that doing so with down-type ones) is called tan β.

The generic sparticle is expected to be heavier than the corresponding particle by an amount $O(M_s)$, though the mass ordering could get reversed for the top + stop, W + chargino and Z + neutralino systems. By assumption, MSSM has a built-in conservation law : that of the multiplicative quantum number R-parity $R_p \equiv (-)^{3B+L+2s}$, with B = baryon no., L = lepton no. and s = spin, which is positive for particles and negative for sparticles. This implies an absolute stability for the lightest sparticle (LSP : a candidate for cold dark matter in cosmology), usually taken to be the lowest-mass neutralino χ_1^0. The LSP, being extremely weakly interacting, escapes through the detectors without leaving any visible trace. The production of sparticle pairs in collider experiments and the consequent decay of each of them is characterized by missing transverse energy E_T signatures. One other consequence of R_p-conservation is the prevention of catastrophic proton decay processes such as $p \to e^+ \pi^0$ which could otherwise proceed with lifetimes $\sim 10^{-8}$ s, instead of $>10^{32}$ yrs as dictated by experiment.

The Lagrangian density of MSSM contains the supersymmetrized minimal extension of that for SM plus the most generally allowed soft supersymmetry breaking (SSB) terms

$$\mathcal{L}^{\text{SM}} \to \mathcal{L}^{\text{MSSM}} = \mathcal{L}^{\text{SSM}} + \mathcal{L}^{\text{SSB}}$$

An attractive feature of MSSM, following from the above, is that couplings among particles and sparticles are simply related by supersymmetry. Some of the vertices, related in these way, are shown below in Figure 4. Note that, in any vertex, sparticles always appear in pairs owing to the constraint of R_p-conservation.

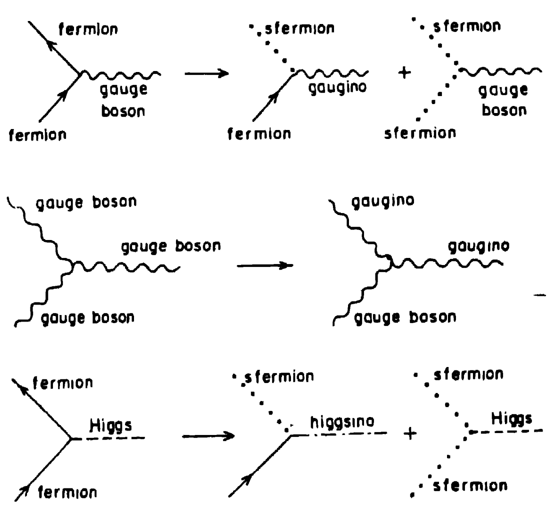


Figure 4. MSMM vertices generated by supersymmetry from those of SM.

More quantitatively, the superfield content of the model in the matter sector, written in a transparent notation ($i = 1, 2, 3$ is a generation index), is :

$$Q_L^i = \begin{pmatrix} U_L^i \\ D_L^i \end{pmatrix}, L_L^i = \begin{pmatrix} N_l^i \\ E_L \end{pmatrix}, Q_R^i, L_R^i, H_1 = \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix}, H_2 \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix}. \tag{2}$$

The corresponding superpotential (with R-parity assumed conserved) is :

$$W = \lambda_U^{ij} Q_L^i . H_2 U_R^j + \lambda_D^{ij} Q_L^i . H_1 D_R^j + \lambda_E^{ij} L_L^i . H_1 E_R^j + \mu H_1 . H_2, \tag{3}$$

with λ's as Yukawa couplings.

The scalar potential can be derived from (3). Writing ϕ_j for a generic scalar field and incorporating the soft supersymmetry breaking terms, we have

$$V = \sum_j \left| \frac{\partial W}{\partial \phi_j} \right| + D-\text{terms} + \sum_{i,j} m_{ij}^2 \phi_i \phi_j + \{ A_U \lambda_U \tilde{q}_L . h_2 \tilde{u}_R$$

$$+ A_D \lambda_D \tilde{q}_L . h_1 \tilde{d}_R + A_E \lambda_E \tilde{e}_L . h_1 \tilde{e}_R + B\mu h_1 . h_2 + H.C. \}. \tag{4}$$

In [4] the third RHS term includes $m_1^2 h_1^+ h_1 + m_2^2 h_2^+ h_2$ with a vanishing m_{12} where $h_{1,2}$ refers to the scalar component of the superfield $H_{1,2}$. Also, $v_{1,2} = \langle h_{1,2}^0 \rangle$ and $\tan\beta = v_2 / v_1$. The physical fields can be expressed in terms of the superfield components given above. For instance, the field for the lightest neutral scalar is $h = \sqrt{2}(\mathrm{Re}\, h_2^0 - v_2)\cos\alpha$ $-\sqrt{2}(\mathrm{Re}\, h_1^0 - v_1)\sin\alpha$, where α is an angle which enters *via* mixing. The orthogonal heavier combination is $H = \sqrt{2}(\mathrm{Re}\, h_2^0 - v_2)\sin\alpha + \sqrt{2}(\mathrm{Re}\, h_1^0 - v_1)\cos\alpha$ while A equals $\sqrt{2}(\mathrm{Im}\, h_2^0 \cos\beta - \mathrm{Im}\, h_1^0 \sin\beta)$. The partners of the CKM matrices in the scalar sector are assumed to possess safety properties which suppress dangerous flavor-changing neutral current processes that could emerge from [3].

At the tree level itself one has several mass relations.

$$m_\pm^2 = m_A^2 + M_W^2, \tag{5a}$$

$$m_h^2 \le M_Z^2 \le m_H^2, \tag{5b}$$

$$\frac{m_h}{|\cos 2\beta|} < m_A < m_H, \tag{5c}$$

$$\mu^2 = (\cos 2\beta)^{-1} \left(m_2^2 \sin^2\beta - m_1^2 \cos^2\beta \right) - \frac{1}{2} M_Z^2, \tag{5d}$$

$$3B\mu = \left(m_1^2 - m_2^2 \right) \tan 2\beta + M_Z^2 \sin 2\beta, \tag{5e}$$

$$m_A^2 = m_1^2 + m_2^2 + 2\mu^2. \tag{5f}$$

On including 1-loop quantum corrections in the leading log approximation, the upper bound on the squared mass of h reads ($\tilde{t}_{1,2}$ are the two physical squarks, assumed to weigh more than the top) [6] :

$$M_h^2 < M_Z^2 \cos^2 2\beta + \frac{3\alpha_{EM}}{2\pi \sin^2 \theta_W} \frac{m_t^4}{M_W^2} \ln \frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2} \simeq (130\ \text{GeV})^2. \qquad (6)$$

This is a "killing prediction" of MSSM.

The renormalization group evolution of the three gauge couplings g_a ($a = 1, 2, 3$) with the energy scale Q are quite different for SM and for MSSM, as shown [7] for $\alpha_a^{-1} \equiv 4\pi g_a^{-2}$ in Figures 5a and 5b. The low energy values of the couplings are now known

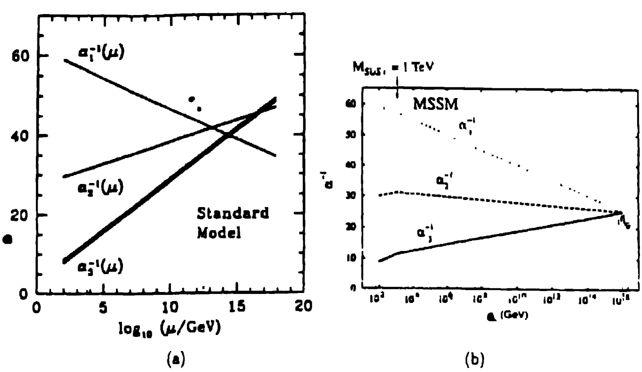


Figure 5. RGE of the gauge couplings in (a) SM and (b) MSSM.

rather accurately and have been used as inputs in these curves. For MSSM the couplings do unify at $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV, while for SM they do not. In Figure 5b M_s has been chosen to be ~1 TeV, but the broad features of the figure do not change when M_s is varied between 100 GeV and 1 TeV. Earlier, when the low energy data were not as precise, SM was compatible with minimal grand unification at ~10^{14} GeV with just a desert in between. Such is no longer the case. This change is illustrated dramatically in the measured values and errors of the sine squared of the Weinberg angle, as shown in Figure 6 for various years starting in 1975. Clearly, grand unified theories, without supersymmetry and basing themselves only on SM at low energies, are ruled out now.

3. CMSSM and the mass spectra of sparticles

Though MSSM is the simplest supersymmetric extension of SM, it introduces 31 new parameters in addition to those of SM. That makes MSSM not very easily testable in terms of predictions that can be pinned down, the predicted upper bound on the lightest Higgs mass [6] being an exception. From a phenomenological standpoint, a more popular version

Plate I

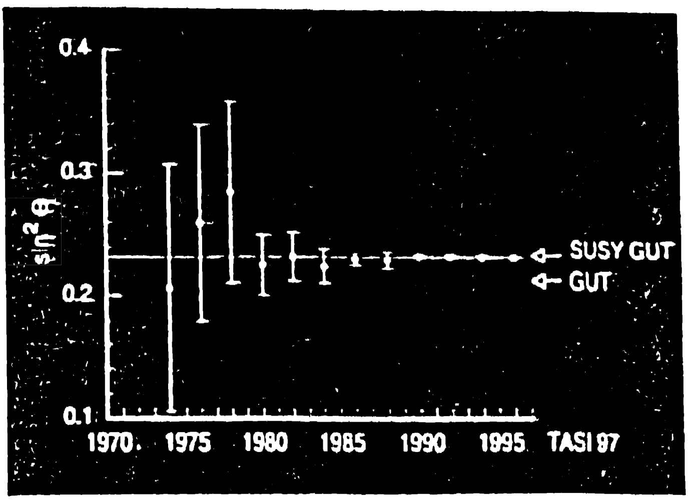


Figure 6. Chronologically progressive reduction of errors in the measurements of $\sin^2\theta_W$.

is the supergravity-constrained [8] MSSM or CMSSM which has only 4 extra parameters plus a sign and hence many dfinitive predictions—especially on the mass spectra of sparticles—that can be tested.

CMSSM has the same Lagrangian density as MSSM. But it is characterized by several simplifying extra assumptions. All of these pertain to boundary conditions (inspired by supergravity theories) imposed on various parameters at the unification scale $M_X \sim 2 \times 10^{16}$ GeV. Specifically, all supersymmetry-breaking scalar (gaugino) masses are assumed to be universal and equal to one mass m_0 ($M_{1/2}$). Squared masses of the Higgs at the unification scale have the additional contribution μ^2 where μ is the supersymmetric Higgsino mass parameter in the MSSM superpotential in [3]. Another assumption is that all supersymmetry-breaking trilinear couplings A_{ijk} in [4] are taken to be equal ($\equiv A_0$). Here m_0 and $M_{1/2}$ are supposedly of the order of the gravitino mass $m_{3/2}$ which sets the scale of M_s. Now m_0, A_0, $M_{1/2}$ and tan β (plus the sign of μ) can be chosen to be the four parameters of CMSSM, or m_A could be traded for one of the first two.

The CMSSM boundary conditions at M_X imply

$$m_1^2(M_X) = m_2^2(M_X) = m_0^2. \tag{7}$$

Turning to gaugino masses M_i (i = nonabelian gauge group index) and considering 1-loop RGE effects, one can write—with α_u as the unified fine structure coupling—

$$M_i(Q) = m_{1/2}\alpha_i(Q)\alpha_u^{-1}(M_X). \tag{8}$$

For the $U(1)_Y$ case, with the standard definition of Y, there is an extra factor of 5/3 in the RHS. It turns out that $M_1(M_Z) \simeq 0.41\ M_{1/2}$ and $M_2(M_Z) \simeq 0.84\ M_{1/2}$ with a mild Q-dependence in $M_{1,2}$. However, the situation is quite different for M_3. The physical on-shell gluino mass $m_{\tilde{g}}$ is given by [9]

$$m_{\tilde{g}} = M_3(Q)\left[1 + \frac{\alpha_S(Q)}{4\pi}\left\{15 - 18\ln\frac{M_3(Q)}{Q} + \sum_q \int_0^1 dx\, x\ln\frac{xm_q^2 + (1-x)m_{\tilde{q}}^2 - x(1-x)M_3^2}{Q^2}\right\}\right] \tag{9}$$

and is independent of Q. For $M_3 \simeq 0.1$ TeV and $m_{\tilde{q}} \simeq 1$ TeV, the difference between $m_{\tilde{g}}$ and M_3 (M_3) can be as much as 30%.

The spectrum of the remaining sparticles can be parametrized, after accounting for renormalization group evolution, as follows [10] :

$$m_{\tilde{e}_R}^2 = m_0^2 + 0.15M_{1/2}^2 - \sin^2\theta_W D; \tag{10a}$$

$$m^2_{\tilde{e}_L} = m_0^2 + 0.52M^2_{1/2} - \left(\frac{1}{2} - \sin^2\theta_W\right)D; \tag{10b}$$

$$m^2_{\tilde{\nu}} = m_0^2 + 0.52M^2_{1/2} + \frac{1}{2}D; \tag{10c}$$

$$m^2_{\tilde{q}\uparrow R} = m_0^2 + (0.07 + C_{\tilde{g}})M^2_{1/2} + \frac{2}{3}\sin^2\theta_W D; \tag{10d}$$

$$m^2_{\tilde{q}\downarrow R} = m_0^2 + (0.02 + C_{\tilde{g}})M^2_{1/2} - \frac{1}{3}\sin^2\theta_W D; \tag{10e}$$

$$m^2_{\tilde{q}\uparrow L} = m_0^2 + (0.47 + C_{\tilde{g}})M^2_{1/2} + \left(\frac{1}{2} - \frac{2}{3}\sin^2\theta_W\right)D \tag{10f}$$

$$m^2_{\tilde{q}\downarrow L} = m_0^2 + (0.47 + C_{\tilde{g}})M^2_{1/2} - \left(\frac{1}{2} - \frac{2}{3}\sin^2\theta_W\right)D. \tag{10g}$$

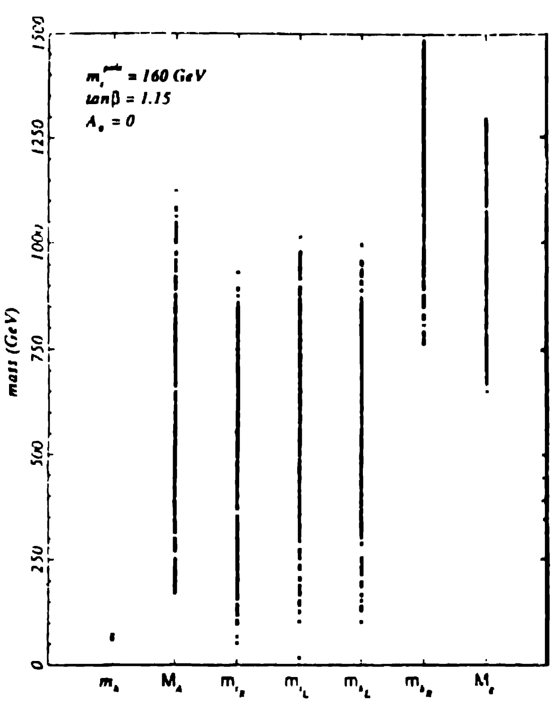


Figure 7. Ranges of some sparticle masses.

Here $C_{\tilde{g}} = \frac{8}{9}[\alpha_S^2(m_{\tilde{q}})/\alpha_S^2(M_X) - 1]$ and $D = M_Z^2\cos^2\beta$ while we have $l = e, \mu$, $q\uparrow = u, c$ and $q\downarrow = d, s, b$. For stops and staus, considerable left-right mixing is anticipated. The corresponding mass-squared matrices are given by

$$m^2_{\tilde{t}} = \begin{pmatrix} m^2_{\tilde{q}\uparrow L} + m_t^2 + 0.35D & -m_t(A_t + \mu\cot\beta) \\ -m_t(A_t + \mu\cot\beta) & m^2_{\tilde{q}\uparrow R} + m_t^2 + 0.16D \end{pmatrix} \tag{11a}$$

$$m^2_{\tilde{\tau}} \qquad \begin{matrix} {}^{2}_{\tilde{l}_L} + m^2_{\tau} - 0.27D & -m_{\tau}(A_{\tau} + \mu \tan\beta) \\ -m_{\tau}(A_{\tau} + \mu \tan\beta) & m^2_{\tilde{l}} + m^2_{\tau} - 0.23D \end{matrix} \Bigg) \tag{11b}$$

We should remark here that arguments exist [11] why tan β should lie between 1 and m_t / m_b.

A sample scatter plot of the ranges [12] of some characteristic masses in the model—showing the extent of variation in the parameter space—is shown in Figure 7. One should also mention that five squarks (*i.e.* all except the stop) need to be taken as nearly mass-degenerate in order to avoid an unacceptable FCNC-induced $K^0 - \bar{K}^0$ mixing. This could be a problem in Figure 7 [12] which has a rather large $\tilde{b}_L - \tilde{b}_R$ mass splitting. A similar argument vis-a-vis the FCNC-induced $\mu \to e\gamma$ decay requires the near mass-degeneracy of all sleptons except $\tilde{\tau}$.

4. Experimental constraints on CMSSM parameters*

In this section, I concentrate on zones in the CMSSM parameter space that can be excluded by use of results from completed or currently running experiments. Some of the constraints, discussed below, involve data from the SLD e^+e^- annihilation experiment at Stanford and the $p\bar{p}$ collision experiments at the Fermilab Tevatron. However, the large majority of them follow from measurements made at the CERN LEP experiments (I will exclude from this talk direct mass limits on squarks and gluinos since those will be covered by D P Roy). The LEP experiments, so far, have an analyzed data sample of more than 20 million Z-peak events at LEP 1 plus nearly 20 pb^{-1} of data in LEP 1.5 at e^+e^- CM energies E_{CM} of 130, 136 and 140 GeV and also about 50 pb^{-1} of data at E_{CM} = 161 GeV.

Let us first state some results in the slepton sector. Sleptons, if accessible in energy, can be pair-produced at LEP. Their characteristic decays with E_T signatures have been looked for. For the right selectron $\tilde{e}_R$, the lower mass bound [13] is $m_{\tilde{e}_R}$ > 75 GeV with the assumption that the mass difference between $\tilde{e}_R$ and the LSP $\tilde{\chi}^0_1$ exceeds 35 GeV. The latter caveat is necessary in the light of the $\tilde{e}_R$ −decay signatures which have been sought in obtaining this bound. For instance, if this mass-difference is taken to exceed only 3 GeV, the said lower mass-bound reduces to 58 GeV. For smuons and staus, the lower mass bounds, with the former assumption, the lower mass bounds (with the former condition) are somewhat weaker, being 55 GeV and 50 GeV respectively, since they get pair-produced only by s-channel processes whereas selectron pair-production has both s- and t-channel contributions. If all sleptons are mass-degenerate and $\tilde{\chi}^0_1$ weighs less than 30 GeV, then the slepton lower mass-bound is 76 GeV.

*This is as of the summer of 1997.

We turn next to light stops. The physical candidates are $\tilde{t}_{1,2}$ with

$$\tilde{t}_1 = \cos\theta \tilde{t}_L + \sin\theta \tilde{t}_R, \tag{12a}$$

$$\tilde{t}_2 = -\sin\theta \tilde{t}_L + \cos\theta \tilde{t}_R \tag{12b}$$

and $\tilde{t}_1$ being lighter. The search process looks for the production $e^+e^- \to \tilde{t}_1\tilde{t}_1^*$, followed by the decay $\tilde{t}_1 \to c\tilde{\chi}_1^0$, so that the final configuration $c\bar{c}\not{E}_T$. The exclusion zones on the $\tilde{\chi}_1^0 - \tilde{t}_1$ mass plot are shown [14] in Figure 8 for $\theta = 0$ and $\theta = \pi/2$ along with the regions excluded by previous LEP 1 and $D\phi$ experiments.

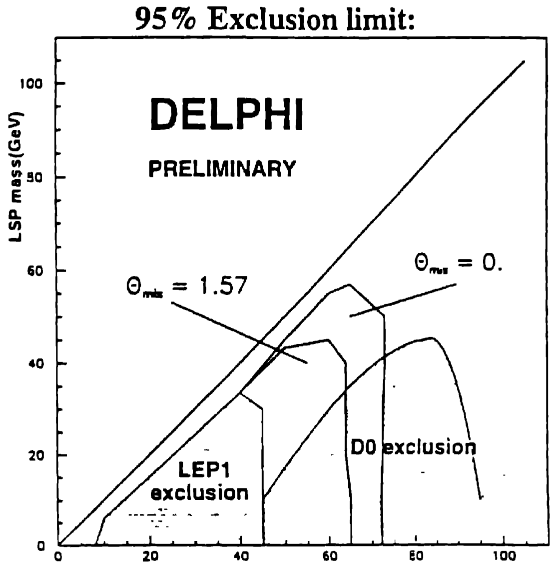


Figure 8. Exclusion zones on the $\tilde{\chi}_1^0 - \tilde{t}_1$ mass plot for extreme values of θ

Coming to the gaugino-higgsino sector now, let us talk spefifically about charginos and neutralinos. Exclusion zones [15] in various mass plots, *i.e.* $\tilde{\chi}_1^0$ *vs* $\tilde{l}_R$ ($\tilde{l} = \tilde{e} + \tilde{\mu} + \tilde{\tau}$), $\tilde{\chi}_1^0$ *vs* $\tan\beta$ and $\tilde{\chi}_1^+$ *vs* $\tilde{\nu}_R$ are shown in Figures 9(a–c) with labels specifying input assumptions. The chargino has been taken to decay by $\tilde{\chi}^\pm \to \tilde{\chi}^0 l^\pm(\tilde{\nu}_l, \tilde{\nu}_l^*)_R$. A distinction has been made between $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^0$ by assuming the former to decay through the process $\tilde{\chi}_2^0 \to \tilde{\chi}_1^0 l^+ l^-$. Assuming that the m_0 parameter is large (> 500 GeV) and that the $M_{1/2}$ parameter is bounded from above by 1 TeV, the following lower mass bounds have been obtained [15] : $m_{\tilde{\chi}_1^0} > 24.6$ MeV, $m_{\tilde{\chi}_2^0} > 32.2$ GeV, $m_{\tilde{}} > 91.1$ GeV, $m > 103.7$ GeV, $m_{\tilde{\chi}_1^\pm} > 73.6$ GeV and $m_{\tilde{\chi}_2^\pm} > 96.2$ GeV.

Coming finally to Higgs scalars, the lightest CP-even supersymmetric Higgs h as well as the CP-odd A have been searched for in the Bjorken process $e^+e^- \to Z \to hZ^* \to b\bar{b}l(q)\bar{l}(\bar{q})$, while both have been sought in LEP 1.5 and LEP 2 in the associated processes $e^+e^- \to Z^* \to Zh \to l(q)\bar{l}(\bar{q})b\bar{b}$ and $e^+e^- \to Z^* \to ZA \to b\bar{b}b\bar{b}, b\bar{b}\tau\bar{\tau}$. The current lower limits are $m_h > 62.5$ GeV, $m_A > 62.5$ GeV for all values of tan β. Stronger lower limits are available for specific assumed values of tan β. In particular, the exclusion zone in the tan β vs m_h plane is shown in Figure 10.

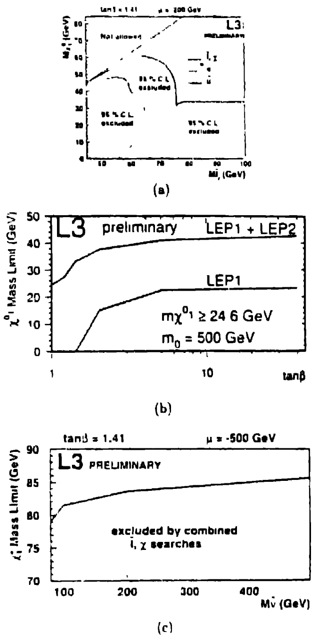


Figure 9. Exclusion zones in various mass-plots : (a) $\tilde{\chi}^0_1$ vs $\tilde{l}_R$, (b) $\tilde{\chi}^0_1$ vs tanβ and (c) $\tilde{\chi}^+_1$ vs $\tilde{\nu}_R$.

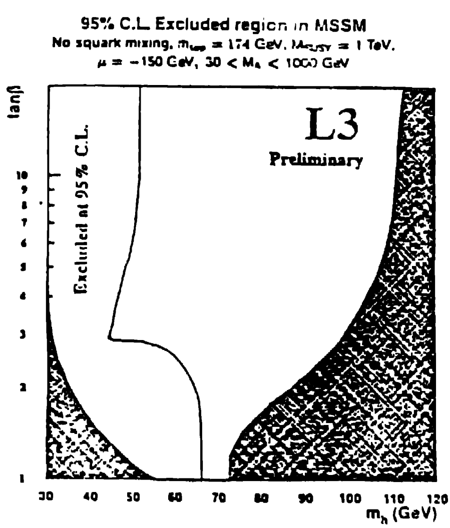


Figure 10. Exclusion zone in the tan β vs m_h plane.

Returning to the parameters of CMSSM, one can choose five independent parameters (m_0, $m_{1/2}$, A_0, tan β and μ). This is tantamount to covering all sfermions (but not

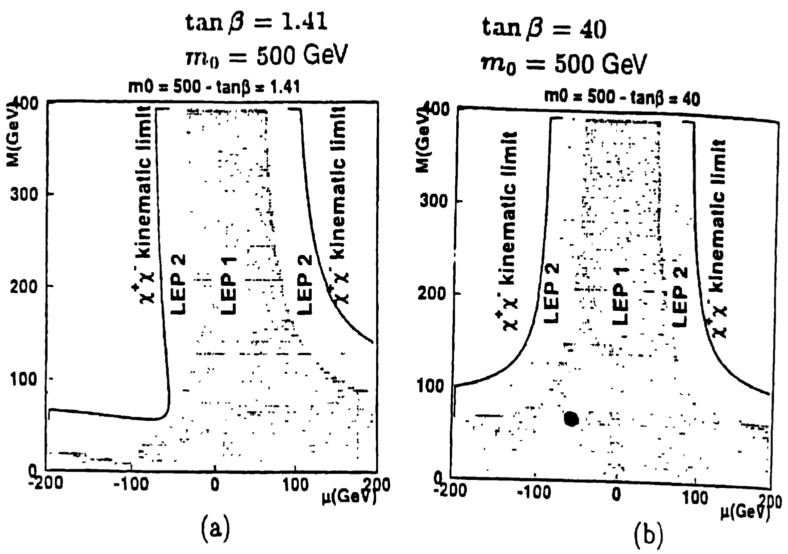
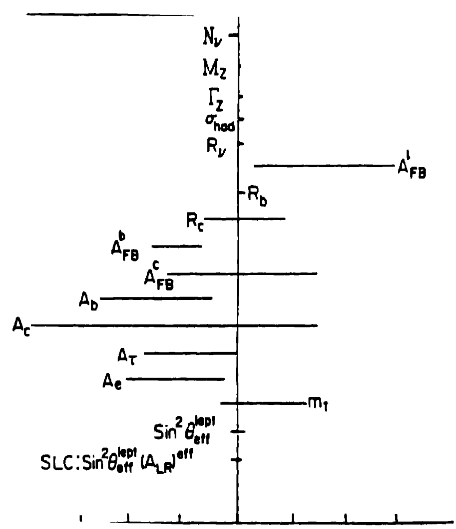


Figure 11. Exclusion zone in the $M_{1/2} - \mu$ plane for $m_0 = 500$ GeV and (a) tan $\beta = 1.41$, (b) tan $\beta = 40$

the Higgs) with the assumption of a universal scalar mass at a high scale. In such a case the exclusion zone, available from the currently analyzed data, in the $M_{1/2} - \mu$ plot is shown

Table 3. Pull-plot, similar to Table 1, for CMSSM



(for m_0 = 500 GeV and tan β = 1.41 or 40) in Figure 11a and 11b. Furthermore, we can compare SM and MSSM fits to the data. The SM fit of Table 1 may be compared with a corresponding "pull-plot" in the CMSSM case shown in Table 3 for tan β = 1.6. The ratio χ^2/(degree of freedom) now is 16.1/12, so that one cannot say that CMSSM is doing significantly better than SM.

5. Conclusions

We can summarize as follows.

(i) Stability considerations of the SM Higgs provide the strongest motivation for near-weak-scale supersymmetry.

(ii) The nature of explicit soft supersymmetry-breaking terms in the low-energy effective Lagrangian is sensitive to input assumptions about high-scale boundary conditions.

(iii) CMSSM, a well-posed theoretical model, is open to challenge from immediate as well as forthcoming experiments.

(iv) The parameter space of CMSSM is getting increasingly restricted as more and more data pour in.

(v) There is a distinct possibility that supersymmetry in nature is decoupled with all sparticles lying near or above 1 TeV.

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References

[1] *LEP Collaboration Report* **CERN-PPE/96-183**

[2] M Schmelling *Proc. 28th Intl Conf. High Energy Physics* (Warsaw, 1996) p 91

[3] Yu A Golfand and E P Likhtmann *JETP Lett* **13** 323 (1971); D V Volkov and V P Akulov *Phys Lett* **46B** 109 (1973); J Wess and B Zumino *Nucl. Phys* **B70** 39 (1974), A Salam and J Strathdee *Nucl Phys.* **B80** 317 (1974)

[4] H E Haber *Proc. Physics from Planck Scale to Electroweak Scale* (Warsaw, 1994) p 49

[5] H E Haber and G L Kane *Phys. Rep.* **117** 75 (1985)

[6] H E Haber and R Hempfling *Phys. Rev. Lett.* **66** 1815 (1991); J Ellis, G Ridolfi and F Zwirner *Phys. Lett.* **B257** 83 (1991); Y Okala, M Yamaguchi and T Yanagida *Phys Lett.* **B262** 54 (1991)

[7] U Amaldi *et al. Phys. Rev.* **D36** 1385 (1987)

[8] For reviews, see H P Nilles *Phys. Rep.* **110** 1 (1984); P Nath, R Arnowitt and A Chamseddine *Applied N=1 Supergravity* (Singapore : World Scientific) (1984); M Drees and S P Martin in *Electroweak Symmetry Breaking and New Physics at the TeV Scale* eds T Barklow, S Dawson, H Haber and J Siegrist (Singapore : World Scientific)

[9] S P Martin and M T Vaughn *Phys. Lett.* **B318** 331 (1993)

[10] M Dress and S P Martin *Ref.* [8]

[11] L E Ibañez and G Ross in *Perspectives on Higgs Physics* ed G L Kane (Singapore : World Scientific) p 229

[12] J L Feng, N Polonski and S Thomas *Phys. Lett* **B370** 95 (1996)

[13] B Barate *et al Aleph Collaboration Report* CERN-PPE/97-056

[14] A De Min *Physics from the Planck Scale to the Electroweak Scale* (*Delphi Collaboration Talk* Given at the 3rd Warsaw Workshop)

[15] S Banerjee *Private Communication* (from the *L3 Collaboration*)